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GRB 080916C: ON THE RADIATION ORIGIN OF THE PROMPT EMISSION FROM KEV/MEV TO GEV

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ABSTRACT

Fermi observations of high-energy gamma-ray emission from GRB 080916C shows that its spectrum is consistent with the Band function from MeV to tens of GeV. Assuming one single emission mechanism dominates in the whole energy range, we show that this spectrum is consistent with synchrotron origin by shock-accelerated electrons. The simple electron inverse-Compton model and the hadronic model are found to be less viable. In the synchrotron scenario, the synchrotron self-Compton scattering is likely to be in the Klein-Nishina regime and therefore the resulting high-energy emission is subdominant, even though the magnetic field energy density is lower than that in relativistic electrons. The Klein-Nishina inverse-Compton cooling may also affect the low-energy electron number distribution and hence results in a low-energy synchrotron photon spectrum $n(\nu)\propto\nu^{-1}$ below the peak energy. Under the framework of the electron synchrotron interpretation, we constrain the shock microphysical parameters and derive a lower limit of the upstream magnetic fields. The detection of synchrotron emission extending to about 70 GeV in the source frame in GRB 080916C favors the Bohm diffusive shock acceleration if the bulk Lorentz factor of the relativistic outflow is not significantly greater than thousands.

Subject headings: gamma rays: bursts

1. INTRODUCTION

It was recently reported that Fermi satellite has detected strong >100MeV emission from a very energetic long-duration burst GRB080916C (Abdo et al. 2009). At a redshift of $z = 4.35\pm0.15$ (Grenier et al. 2009), the burst is the most energetic one ever, with an isotropic gamma-ray energy $E_\gamma\simeq 8\times10^{54}$ergs, which is released over a duration of $T_{90}\simeq 60$s. Equally remarkably, more than ten photons with energy above GeV are detected, with the highest energy one at 13 GeV (in the observer frame). The spectra of all five time intervals, designated as times a-e in the light curves of GRB080916C (Abdo et al. 2009), are well fit by the empirical Band function (Band et al. 1993), which smoothly joins low- and high-energy power laws. The high-energy power law extends to GeV energies, without any additional spectral component visible. The peak energy of the spectra during these intervals is around $\varepsilon_p\simeq 400\mathrm{KeV}-1\mathrm{MeV}$. Except during the first time interval, the low-energy and high-energy photon spectral indices of the prompt emission are constant and consistent with $\alpha\simeq-1.0$ and $\beta\simeq-2.2$ respectively. With its high temporal resolution, INTEGRAL detected the temporal variability of the KeV/MeV emission on time scales as short as 100 ms with high statistical significance (Greiner et al. 2009). So the variability timescale in the local source frame is $t_v\lesssim 100\mathrm{ms}/(1+z)=20\mathrm{ms}$.

The nonthermal synchrotron radiation by electrons has been suggested to be a possible mechanism for the 10 KeV-MeV emission (see Mészáros 2006 and Zhang 2007 for recent reviews), but one famous problem remains so far, i.e. the low-energy photon spectral index α is incompatible with the index $-3/2$ that is expected from fast-cooling electrons (e.g. Preece et al. 1996; Ghisellini et al. 2000). Electron inverse Compton emission has been a competitive mechanism (e.g. Panaitescu & Mészáros 2000). The fact that one single spectral component fits the spectrum of the prompt emission from 10 KeV to GeV in GRB080916C suggests that one emission mechanism dominates in this whole energy range[6]. In this *Letter*, we study the constraint that this puts on the emission mechanism. Abdo et al. (2009) mentioned as one of the possibilities that the delay of high-energy gamma-ray emission relative to low-energy emission in GRB080916C could be a result of longer acceleration time needed for higher energy protons or nuclei in hadronic emission models. In accordance with this, we also study whether such hadronic models could be a possible mechanism that produces the KeV/MeV to GeV emission in GRB080916C.

2. THE SYNCHROTRON MODEL AND PARAMETER CONSTRAINTS

Assuming that in GRB shocks, fractions of ϵ_B and ϵ_e of the shock internal energy are converted into the energy in the magnetic field and electrons, respectively. To ensure a high radiation efficiency for the prompt emission, it is usually assumed that the electrons are rapidly cooling, so the energy density in gamma-ray emission U_γ is equal to the electron energy density U_e. The magnetic field is given by

$$\frac{B^2}{8\pi}=\left(\frac{\epsilon_B}{\epsilon_e}\right)U_\gamma=\left(\frac{\epsilon_B}{\epsilon_e}\right)\frac{L_\gamma}{4\pi R^2c\Gamma^2}, \qquad (1)$$

where R is the radius of the shock, L_γ is the luminosity in gamma-ray emission and Γ is the bulk Lorentz factor. The detection of GeV photons suggest that the emission region has a bulk Loentz factor $\Gamma\gtrsim10^3$(Grenier et al. 2009; Abdo et al. 2009)[7]. From the casuality constraint, the emission radius is $R=\Gamma^2ct_v$. In the synchrotron model for the 10 KeV-GeV

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[6] Dropping the assumption of the same origin for both MeV and high energy emission, Li (2008) explained the high energy emission in GRB080916C as delayed, "residual" emission from subsequent collisions at larger and larger radii in the baryonic outflow.

[7] The limit on the bulk Lorentz factor would be more stringent when we have an IC TeV component whose flux is above the synchrotron extension,

emission, by use of $\varepsilon_p = h\nu_p = \phi_\nu \frac{3hqB}{4\pi m_e c}\gamma_m^2\Gamma$, one can derive the Lorentz factor of electrons that radiate at the GRB peak energy ε_p,

$$\begin{aligned}\gamma_m &= (\tfrac{4\pi m_e c\varepsilon_p}{3\phi_\nu hq})^{1/2}(\tfrac{\epsilon_e}{\epsilon_B})^{1/4}(\tfrac{2L_\gamma}{R^2 c})^{-1/4}\varepsilon_p^{1/2}\\ &= 2.5\times10^3(\tfrac{\epsilon_e}{\epsilon_B})^{1/4}L_{\gamma,54}^{-1/4}\Gamma_3 t_{v,-2}^{1/2}(\tfrac{\varepsilon_p}{2\mathrm{MeV}})^{1/2},\end{aligned} \tag{2}$$

where $\phi_\nu \simeq 0.5$ is the coefficient defined in Wijers & Galama (1999) and q is the electron charge. Define $\gamma_{\rm T}$ as the Lorentz factor of electrons below which the scattering with peak energy photons are in the Thomson scattering regime, i.e.

$$\gamma_{\rm T} = \Gamma m_e c^2/\varepsilon_p = 250\Gamma_3(\varepsilon_p/2\mathrm{MeV})^{-1}. \tag{3}$$

Unless $\epsilon_e \lesssim 10^{-4}\epsilon_B$, which is unreasonable in terms of the burst energetics, the IC scattering between γ_m electrons and the bulk of the gamma-ray emission should be in the Klein-Nishina (KN) regime.

The KN Compton cooling of electrons may affect the low-energy electron distribution at $\gamma_{\rm T} < \gamma < \gamma_m$ and hence affect the low-energy spectral slope below ε_p (e.g. Rees 1967; Derishev et al. 2003), as we show below. Consider a population of electrons injected into a box with magnetic field B in a power law form $N(\gamma) \propto \gamma^{-p}$ for $\gamma \geq \gamma_m$. These electrons will cool down rapidly through synchrotron and/or IC radiation[8]. As the electron cools, its energy changes as γ, so we have

$$\nu F_\nu\left[1 + k(\gamma)U_\gamma/U_B\right] \propto \gamma \tag{4}$$

for $\gamma < \gamma_m$, where ν is the synchrotron frequency of γ-electrons and $k(\gamma)$ accounts for the reduction of the effective photon energy density for IC scattering of the γ-electrons due to the KN effect. Define $h\nu_{KN} = \Gamma m_e c^2/\gamma$ as the critical energy of the photons with which the IC scattering of γ-electrons is just in the KN regime. For a low-energy photon spectrum of the form $\nu F_\nu \propto \nu^\delta$ $(\nu < \nu_p)$, we have

$$k(\gamma) \simeq \frac{U_\gamma(\nu<\nu_{KN})}{U_\gamma} = (\frac{\nu_{KN}}{\nu_p})^\delta = (\frac{\gamma}{\gamma_{\rm T}})^{-\delta} \tag{5}$$

for $\gamma_{\rm T} < \gamma < \gamma_m$ and $k \simeq 1$ for $\gamma < \gamma_{\rm T}$. In the case of $\gamma_{\rm T} < \gamma < \gamma_m$, if $kU_\gamma/U_B \gg 1$, i.e. the IC cooling is still dominant even though the scatterings are in the KN regime, one can obtain $\nu F_\nu \propto \gamma^{\delta+1} \propto \nu^{(\delta+1)/2}$, where we have used $\nu \propto \gamma^2$ in the last step. Equating this synchrotron spectral index with the initially assumed spectral index $F_\nu \propto \nu^{\delta-1}$, one can derive

$$\delta = 1, F_\nu \propto \nu^0. \tag{6}$$

This can explain the observed low-energy photon index of $\alpha = -1.0$ in GRB080916C if the condition, $U_B \lesssim k(\gamma)U_\gamma$ for $\gamma_{\rm T} < \gamma < \gamma_m$, is satisfied.

On the other hand, the high-energy spectrum of GRB080916C above ε_p is consistent with the synchrotron spectrum produced by fast-cooling electrons above γ_m, i.e. $F_\nu \propto \nu^{-p/2}$ with $p = -2(1+\beta_2) = 2.4$. The dominance of synchrotron cooling above ε_p implies that $U_B \gtrsim k(\gamma)U_\gamma$ for electrons with Lorentz factor $\gamma \gtrsim \gamma_m$. Therefore, we find that, at $\gamma = \gamma_m$, $U_B \simeq kU_\gamma$. Since $U_\gamma \simeq U_e$ for fast-cooling electrons, the requirement $U_B \simeq k(\gamma_m)U_\gamma$ translates into $\epsilon_e/\epsilon_B \simeq \gamma_m\varepsilon_p/\Gamma m_e c^2$, which gives

$$\frac{\epsilon_e}{\epsilon_B} \simeq 20L_{\gamma,54}^{-1/3}t_{v,-2}^{2/3}(\frac{\varepsilon_p}{2\mathrm{MeV}})^2. \tag{7}$$

Note that the transition region between the two asymptotic power-laws at low and high energy ends in the Band function is rather wide, so the above requirement, $U_B \simeq k(\gamma_m)U_\gamma$, should be regarded as an order of magnitude of estimate. In addition, this requirement applies only to large ε_p bursts, because for low ε_p bursts, the IC scattering may be no longer in the Kelin-Nishina regime. A signature that high ε_p bursts have α preferentially close to -1 can be seen in the analysis of Preece et al. (1996). For some low ε_p bursts that have $\alpha \simeq -1$, some other mechanisms may be at work.

For electrons with $\gamma \lesssim \gamma_{\rm T}$, the IC scatterings with peak energy photons are in the Thomson scattering regime, so $k(\gamma) = 1$ and $N(\gamma) \propto \gamma^{-2}$, leading to a conventional fast-cooling photon spectrum of $F_\nu \propto \nu^{-1/2}$. Observations show a single power law spectrum $F_\nu \propto \nu^0$ from 10 keV to $\sim$ MeV in GRB080916C, implying that $\gamma_m/\gamma_{\rm T} \gtrsim 10$, and one can therefore obtain a constraint

$$\frac{\epsilon_e}{\epsilon_B} \gtrsim 1L_{\gamma,54}t_{v,-2}^{-2}(\frac{\varepsilon_p}{2\mathrm{MeV}})^{-6} \tag{8}$$

Due to that the IC scatterings between γ_m electrons and the peak energy photons with energy ε_p are in the KN regime, the IC emission peaks at

$$\begin{aligned}h\nu_p^{IC} &= \Gamma\gamma_m m_e c^2\\ &= 1(\tfrac{\epsilon_e}{\epsilon_B})^{1/4}L_{\gamma,54}^{-1/4}\Gamma_3^2 t_{v,-2}^{1/2}(\tfrac{\varepsilon_p}{2\mathrm{MeV}})^{1/2}\mathrm{TeV},\end{aligned} \tag{9}$$

with a flux $\nu F_\nu^{IC}(\varepsilon_\gamma = h\nu_p^{IC}) = Y(\gamma_m)\nu_p F_{\nu_p} \simeq \nu_p F_{\nu_p}$, where Y is the Compton parameter. For a flat synchrotron spectrum with $\beta \simeq -2$ above ν_p, it is natural that the IC component is not seen at high energies since $\nu F_\nu^{syn}(\varepsilon_h = 70\mathrm{GeV}) \gtrsim \nu F_\nu^{IC}(\varepsilon_h = 70\mathrm{GeV})$ for GRB080916C.

In the above, we have not assumed any model for the dissipation mechanism of the shocks. In the popular internal shock scenario, the typical Lorentz factors of the shocked electrons is $\gamma_m = \epsilon_e(m_p/m_e)\Gamma_{sh}$, where Γ_{sh} is the shock Lorentz factor, which is equal to the relative Lorentz factor of the two colliding shells. For GRB080916C, we have obtained a constraint $\gamma_m \simeq 5\times10^3\Gamma_3 L_{\gamma,54}^{-1/3}t_{v,-2}^{2/3}(\frac{\varepsilon_p}{2\mathrm{MeV}})$. So if internal shock applies to GRB080916C, we would need a large relative Lorentz factor, $\Gamma_{sh} \simeq 8(3\epsilon_e)^{-1}\Gamma_3 L_{\gamma,54}^{-1/3}t_{v,-2}^{2/3}(\frac{\varepsilon_p}{2\mathrm{MeV}})$. This could be caused by the interaction among the shells whose Lorentz factors have a large contrast (Yu et al. 2009). Of course, the shock could also arise from the magnetic reconnection or turbulence (e.g. Thompson 1994; Mészáros & Rees 1997; Lyutikov & Blandford 2003; Narayan & Kumar 2008; Zhang & Pe'er 2009) and we do not have the estimate of the shock Lorentz factor from the first principle.

3. ALTERNATIVE MODELS FOR MEV-10 GEV EMISSION?

3.1. *The one-zone SSC scenario*

Let's explore whether the simple IC scenario (i.e. one zone synchrotron self-Compton (SSC) scenario) can explain the single power-law spectrum from MeV to 10 GeV in GRB080916C. Suppose that the first-order SSC of electrons with energy γ_m produce the peak emission $\varepsilon_p = 2$MeV. Since this IC emission is not hidden by the synchrotron emission, one will expect that the 2nd order IC emission appears at high-energy if the 2nd-order IC peak is located within the observation energy window and that the 2nd-order IC scattering is still in the Thomson scattering regime. The fact that we did not see the 2nd-order IC component implies that $\gamma_m^2\varepsilon_p \gtrsim 70$GeV or $\gamma_m\varepsilon_p \gtrsim \Gamma m_e c^2$, so we have $\gamma_m \gtrsim$

as is the case in the synchrotron scenario. Using the formula in Lithwick & Sari (2001), however, one can find that the limit is only increased by a factor smaller than two.

[8] The synchrotron cooling time of γ_m electrons is shorter than the dynamic time $t_d' = R/\Gamma c = \Gamma t_v$ as long as $\epsilon_e/\epsilon_B \lesssim 10^6\Gamma_3^{-16/3}t_{v,-2}^{-2}(\varepsilon_p/2\mathrm{MeV})^{2/3}L_{\gamma,54}$.

$190(\varepsilon_p/2\mathrm{MeV})^{-1/2}$ or $\gamma_m \gg \gamma_{\rm T} = 250\Gamma_3(\varepsilon_p/2\mathrm{MeV})^{-1}$. Since $\varepsilon_p = h\nu_{syn,p}\gamma_m^2$, one obtains the synchrotron peak frequency at $h\nu_{syn,p} = 55(\gamma_m/190)^{-2}(\varepsilon_p/2\mathrm{MeV})\mathrm{eV}$. Then one can obtain an upper limit of the magnetic filed $B = h\nu_{syn,p}/(\phi_\nu \frac{3q}{4\pi m_e c}\gamma_m^2\Gamma) = 140(\gamma_m/190)^{-4}\Gamma_3^{-1}(\varepsilon_p/2\mathrm{MeV})\mathrm{G}$. With this magnetic field, we can derive an upper limit for ϵ_B/ϵ_e. In the case that 2nd-order IC is still in the Thomson regime, $U_e = (Y+1)U_\gamma$, so from $\epsilon_B/\epsilon_e = U_B/U_e \simeq U_B/(YU_\gamma)$, we obtain $Y(\epsilon_B/\epsilon_e) = 3\times10^{-5}(\gamma_m/190)^{-8}(\varepsilon_p/2\mathrm{MeV})^2 L_{\gamma,54}^{-1}\Gamma_3^4 t_{v,-2}^2$. By use of $Y = (\epsilon_e/\epsilon_B)^{1/3}$ (Kobayashi et al. 2007), one get $Y = 170(\gamma_m/190)^4(\varepsilon_p/2\mathrm{MeV})^{-1}L_{\gamma,54}^{1/2}\Gamma_3^{-2}t_{v,-2}^{-1}$. As $\gamma_m \gtrsim 190$, so the radiation energy in the 2nd-order IC will be $E_{2nd,IC} = YE_\gamma \gtrsim 1.3\times10^{57}$erg, which is too large to be realistic. Such an energy crisis problem has also been found in the case of GRB080319B for the IC scenario of the prompt MeV emission (Piran et al. 2008).

On the other hand, if the 2nd-order IC is already in the deep KN regime (for $\gamma_m \gg \gamma_T = 250\Gamma_3$), $\epsilon_B/\epsilon_e = U_B/U_e = U_B/U_\gamma$. The 2nd-order Compton Y parameter is $Y_{\rm 2nd} = (\epsilon_e/\epsilon_B)^{1/2}(\gamma_m/\gamma_{\rm T})^{-\delta}$ for a spectrum $\nu F_\nu \propto \nu^\delta$ below ε_p (see Eq.5). To get $Y_{\rm 2nd} \lesssim 1$, one needs $\gamma_m \gtrsim (\epsilon_e/\epsilon_B)^{1/2}\gamma_{\rm T} = 250(\epsilon_e/\epsilon_B)^{1/2}\Gamma_3(\varepsilon_p/2\mathrm{MeV})^{-1}$ for $\delta = 1$. So from $h\nu_{syn,p} = 55(\gamma_m/190)^{-2}(\varepsilon_p/2\mathrm{MeV})\mathrm{eV}$, one can obtain $B \lesssim 45(\epsilon_e/\epsilon_B)^{-2}\Gamma_3^{-1}(\varepsilon_p/2\mathrm{MeV})\mathrm{G}$. Combing this upper limit with the equipartition assumption in Eq.(1), one can get $\epsilon_e/\epsilon_B \lesssim 0.02\Gamma_3^{4/3}t_{v,-2}^{2/3}(\varepsilon_p/2\mathrm{MeV})^{2/3}L_{\gamma,54}^{-1/2}$, which is in conflict with the precondition $\epsilon_e/\epsilon_B = Y^2 \gtrsim 1$. This means that significant suppression of the 2nd-order IC emission by KN scattering can not be fulfilled. So we conclude that the the simple one-zone SSC model does not work for the MeV to 10 GeV emission in GRB080916C.

3.2. *The hadronic scenario*

We first study whether the proton synchrotron emission can produce the MeV-10 GeV emission of GRB080916C. The photon spectrum index above ε_p, $\beta = -2.2$, implies that the proton distribution index is $p \simeq 2.4$ for fast-cooling protons or a very steep index $p \simeq 3.4$ for slow-cooling ones. In the proton synchrotron scenario, the observed peak emission at ε_p is produced by protons with a Lorentz factor of $\gamma_p = (\frac{4\pi\varepsilon_p m_p c}{3\phi_\nu qhB\Gamma})^{1/2}$. The synchrotron cooling time of these protons in the comoving frame is $t'_{syn} = 6\pi m_p^3 c/(\sigma_T m_e^2\gamma_p B^2)$. Define that the magnetic field energy density is a fraction of ξ_B of the comoving frame photon energy density, i.e. $U_B = \xi_B U_\gamma$. So the synchrotron radiation efficiency of the γ_p protons is $\eta = \min[1, t'_d/t'_{syn}]$, where $t'_d = R/\Gamma c$ is the dynamic time in the comoving frame, which is equal to the comoving frame variability time, $t'_d = t'_v = \Gamma t_v$. As long as $t'_v/t'_{syn} \lesssim 1$, we have a radiation efficiency for γ_p protons

$$\eta(\gamma_p) = \frac{t'_v}{t'_{syn}} = 3\times10^{-4}\xi_B^{3/4}L_{\gamma,54}^{3/4}t_{v,-2}^{-1/2}\Gamma_3^{-4}(\frac{\varepsilon_p}{2\mathrm{MeV}})^{1/2} \qquad (10)$$

Such a low radiation efficiency implies an unrealistically large energy in protons, a factor of $1/\eta \simeq 3\times10^3\xi_B^{-3/4}$ higher than the energy in gamma-rays. The radiation efficiency is quite low ($\eta \sim 2\times10^{-3}$) even for protons that produce the high-energy gamma-rays of energy $\gtrsim 100$MeV. So there is no room for the proton synchrotron model even in the assumption that the high-energy gamma-ray emission belongs to a different component than the MeV component. If the spectrum above ε_p is interpreted as arising from fast-cooling protons, as required in the case of $2 \lesssim p \lesssim 3$, one would need $\xi_B = U_B/U_\gamma \gtrsim 5\times10^4 L_{\gamma,54}^{-1}t_{v,-2}^{2/3}\Gamma_3^{16/3}(\frac{\varepsilon_p}{2\mathrm{MeV}})^{-2/3}$, which is also unreasonable, as the total energy in the magnetic field is too large to be realistic for a GRB.

Let's also explore the scenario of the secondary emission from hadronic photopion process. Detection of high-energy gamma-rays of energy greater than 10 GeV puts a constraint on the opacity of $\gamma\gamma$ absorption. As the hadronic $p\gamma$ opacity is related with the $\gamma\gamma$ opacity, a higher maximum photon energy, hence a lower $\gamma\gamma$ absorption opacity, would imply a lower hadronic radiation efficiency (e.g. Dermer et al. 2008). It is useful to express the hadronic $p\gamma$ efficiency as a function of the the pair production optical depth $\tau_{\gamma\gamma}$. Following Waxman & Bahcall (1997), the optical depth for pair production of a photon of energy ε_h is $\tau_{\gamma\gamma}(\varepsilon_h) = \frac{R}{\Gamma l_{\gamma\gamma}} = \frac{R}{\Gamma}\frac{\sigma_T}{16}\frac{U_\gamma\varepsilon_h}{\Gamma(m_ec^2)^2}$, where $l_{\gamma\gamma}$ is the mean free path. For the simplicity of calculation, we have assumed a photon spectrum $\beta_2 = -2$ above ε_p, which is a good approximation for GRB080916C. The fraction of energy lost by protons to pions is $f_\pi \simeq \frac{R}{\Gamma}\frac{U_\gamma}{2\varepsilon'_p}\sigma_{p\gamma}\xi_{\rm peak}$ for protons with energy greater than $E_p = 6\times10^{16}\Gamma_3^2(\varepsilon_p/2\mathrm{MeV})^{-1}$eV (Waxman & Bahcall 1997), where $\sigma_{p\gamma} \simeq 5\times10^{-28}\mathrm{cm}^2$ is the cross section of the $p\gamma$ reaction at the Δ resonance and $\xi_{\rm peak} \simeq 0.2$ is the fraction of proton energy loss in one interaction. So the maximum photopion efficiency is

$$\begin{aligned} f_\pi &= 2\times10^{-3}\Gamma_3^2(\tfrac{\varepsilon_h}{70\mathrm{GeV}})^{-1}(\tfrac{\varepsilon_p}{2\mathrm{MeV}})^{-1}\tau_{\gamma\gamma}(\varepsilon_h) \\ &= 2\times10^{-3}(\Gamma/\Gamma_{\rm lim})^{-6}\Gamma_3^2(\tfrac{\varepsilon_h}{70\mathrm{GeV}})^{-1}(\tfrac{\varepsilon_p}{2\mathrm{MeV}})^{-1} \end{aligned} \qquad (11)$$

where $\tau_{\gamma\gamma} \simeq (\Gamma/\Gamma_{\rm lim})^{-6}$ has been used in the last step and $\Gamma_{\rm lim}$ is the minimum bulk Lorentz factor of the outflow required by the transparency of the high energy photon of energy ε_h. As $\tau_{\gamma\gamma}(\varepsilon_h = 70\mathrm{GeV}) \lesssim 1$, we obtain an upper limit of the $p\gamma$ efficiency, i.e. $f_\pi \lesssim 2\times10^{-3}$. If the prompt MeV-10 GeV emission is interpreted as arising from the secondary emission of hadronic process, one would need an unreasonably large energy budget in relativistic protons.

4. THE MAXIMUM SYNCHROTRON PHOTON ENERGY AND ITS IMPLICATIONS

First we calculate the maximum synchrotron photon energy that can be reached in GRB shocks. We assume that the relative motion between the upstream and downstream plasma is only mildly relativistic (such as in internal shocks). For an electron being shock accelerated, the residence time in downstream and upstream regions are respectively, $t'_d = \kappa_d\varepsilon'_{e,d}/qB_dc$ and $t'_u = \kappa_u\varepsilon'_{e,d}/qB_uc$, where ε'_e is the energy of accelerated electrons, B_d and B_u are respectively the magnetic fields in the shock downstream and upstream, and $\kappa_{d,u} \gtrsim 1$ parameterizes the efficiency of shock acceleration, with $\kappa_{d,u} \simeq 1$ corresponding to the fastest shock acceleration– the Bohm diffusive shock acceleration with the scattering mean free path equal to the particle gyroradius. It is generally assumed that the downstream magnetic field is close to the equipartition with the shock internal energy, while the value of upstream magnetic field is less clear. As $B_u \lesssim B_d$, the total acceleration time is dominated by upstream residence time, so $t'_{acc} \simeq \kappa\varepsilon'_e/qB_uc$. The maximum energy of accelerated electrons in each region is determined by equating the residence time with the shorter one of the cooling time and the available dynamic time, i.e. $t'_{acc} = \min\{t'_{cool}, t'_{dyn}\}$. The cooling time in the downstream and upstream are, respectively, $t'_{cool,d,u} = 3m_ec/(4\sigma_T(U_{B_{d,u}}+k(\gamma)U_\gamma))$, where U_{B_d} and

U_{B_u} represent the magnetic field energy density in downstream and upstream respectively, and U_γ is the photon energy density. In downstream region, $U_{B_d} \gg k(\gamma_{M,d})U_\gamma$, so the maximum electron energy is $\gamma_{M,d} = (\frac{6\pi q}{\kappa_d \sigma_T B_d})^{1/2}$. In upstream region, the magnetic field energy density could be lower than $k(\gamma_{M,u})U_\gamma$, and in this case, the maximum electron energy is $\gamma_{M,u} = (\frac{3qB_u}{4\kappa_u\sigma_T k(\gamma_{M,u})U_\gamma)})^{1/2}$, where $k(\gamma_{M,u}) = (\gamma_{\rm T}/\gamma_{M,u})^{1/2}$. So $\gamma_{M,u} = (\frac{3qB_u}{4\kappa_u\sigma_T U_\gamma \gamma_{\rm T}^{1/2}})^{2/3}$. As $B_d \gtrsim B_u$, the electrons radiate more efficiently in the downstream and therefore the relevant maximum Lorentz factor with the observed radiation is $\gamma_M = \min[\gamma_{M,u}, \gamma_{M,d}]$. Depending on which of $\gamma_{M,u}$ and $\gamma_{M,d}$ is larger, we divide the discussion into two cases:

i)The $\gamma_{M,d} \lesssim \gamma_{M,u}$ case. The maximum synchrotron photon energy is

$$h\nu_{syn,M} = 0.2294 \frac{3qB_d}{4\pi m_e c}\gamma_{M,d}^2 \Gamma = 55\left(\frac{1}{\kappa_d}\right)\Gamma_3 \mathrm{GeV}, \quad (12)$$

which is only dependent of the bulk Lorentz factor Γ of the relativistic outflow (0.2294 is the coefficient quoted from Wijers & Galama 1999). If the observed highest energy photon is produced by synchrotron radiation, from $h\nu_{syn,M} \gtrsim \varepsilon_h$, we obtain

$$\kappa_d \lesssim 0.8\Gamma_3 \left(\frac{\varepsilon_h}{70\mathrm{GeV}}\right)^{-1}, \quad (13)$$

which favors the Bohm diffusive acceleration if the bulk Lorentz factor $\Gamma_3 \lesssim a\,few$.

From the precondition, $\gamma_{M,d} \lesssim \gamma_{M,u}$, we obtain a lower limit of the upstream magnetic field in this case,

$$\begin{aligned} B_u &\gtrsim \frac{4\kappa_u\sigma_T U_\gamma \gamma_{\rm T}^{1/2}}{3q}\left(\frac{6\pi q}{\kappa_d \sigma_T B_d}\right)^{3/4} \\ &= 500(\tfrac{\epsilon_e}{\epsilon_B})^{3/8}\kappa_u \kappa_d^{-3/4} L_{\gamma,54}^{5/8}\Gamma_3^{-13/4}(\varepsilon_p/2\mathrm{MeV})^{-1/2} t_{v,-2}^{-5/4}\,\mathrm{G} \end{aligned} \quad (14)$$

ii)For the case $\gamma_{M,u} \lesssim \gamma_{M,d}$, the maximum synchrotron photon energy is $h\nu_{syn,M}(\gamma_{M,u}) = 0.2294\frac{3qhB_d}{4\pi m_e c}\gamma_{M,u}^2\Gamma$. So from $h\nu_{syn,M}(\gamma_{M,u}) \gtrsim \varepsilon_h$, we obtain a lower limit of the upstream magnetic field,

$$\begin{aligned} B_u &\gtrsim \left(\frac{4\pi m_e c E_h}{0.2294\times 3qB_d\Gamma}\right)^{3/4}\frac{4\kappa_u\sigma_T U_\gamma \gamma_{\rm T}^{1/2}}{3q} \\ &\simeq 600\kappa_u(\tfrac{\epsilon_e}{\epsilon_B})^{3/8}\left(\tfrac{\varepsilon_h}{70\mathrm{GeV}}\right)^{3/4} L_{\gamma,54}^{5/8}\Gamma_3^{-4}(\tfrac{\varepsilon_p}{2\mathrm{MeV}})^{-1/2} t_{v,-2}^{-5/4}\mathrm{G}. \end{aligned} \quad (15)$$

Combining this lower limit with the precondition $\gamma_{M,u} \lesssim \gamma_{M,d}$, we find that Eq.(15) is applicable only when $\varepsilon_h \lesssim 55\kappa_d^{-1}\Gamma_3$GeV. In both cases, the GRB shells that produce the prompt emission must have a pre-shock magnetic field greater than ~ 500G at a radius of $R \sim 3\times 10^{14}\mathrm{cm}\Gamma_3^2 t_{v,-2}$. If the field lines in the expanding shell are frozen and the width of the shell is constant, the components then vary with distance as $B_r \propto r^{-1}$ and $B_\theta \sim B_\phi \sim r^{-2}$. For an initial magnetic filed of $B_0 \sim 10^{15}$G within a volume of radius of $10^6 - 10^7$cm, the above limit is larger than the B_r component, but still within the B_θ or B_ϕ component. Of course, the above limit is also consistent with the hypothesis that the upstream magnetic field is significantly amplified by the particle streaming instability (Bell 2004)[9]. Interestingly, the shock compressed upstream magnetic field, $B \sim 4\Gamma_{sh}B_u \gtrsim 1500$G, is similar to the assumed equipartition magnetic field in downstream (i.e. Eq.1), which means that the field compression due to the shock is enough to explain the downstream magnetic field.

[9] Recently, Li & Waxman (2006) constrained the pre-shock magnetic fields of GRB afterglow shocks by synchrotron X-ray afterglows, which also implies that the pre-shock magnetic fields may be amplified.

5. SUMMARY AND DISCUSSIONS

The single-component spectrum of GRB080916C from MeV to GeV puts useful constraints on the emission mechanism. We found that the synchrotron mechanism from relativistic electrons is consistent with the observed spectrum, while the simple one-zone electron IC and hadronic models are less viable. In the synchrotron interpretation, the SSC emission is found to be in the KN scattering regime and as a consequence, the IC component is not visible at high energies even though the magnetic field energy density is smaller than that in the relativist electrons, i.e. $\epsilon_B < \epsilon_e$, as obtained in our case. We also suggest a scenario in which such a KN IC emission dominated regime can explain the low energy photon spectral index of GRB 080916C.

The delay of high-energy gamma-ray emission relative to the low-energy emission in GRB080916C is still a mystery in the electron synchrotron scenario. It could be due to that the energy distribution slope p of electrons during the first time interval (time a) is rather steep so that the high-energy emission is suppressed or that the emission region has not become transparent for high-energy gamma-rays at early times.

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